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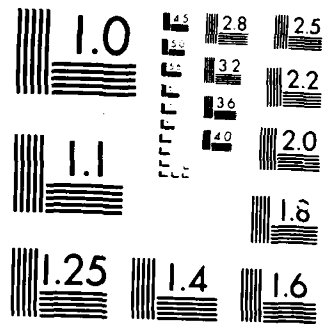

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ABSTRACT

A numerical and analytical study of soliton propagation and generation in a Raman amplifier is performed. Parameters are amplitude and temporal profile of input optical fields, Raman line width and detuning.

Soliton excitations will develop, if a phase shift of nearly 180 degrees is introduced into the Stokes seed beam. In the absence of detuning and for an instantaneous phase shift of exactly 180 degrees the solitons will be stable. In the presence of detuning and for non instantaneous phase shifts of less than 180 degrees the solitons will be unstable and decay for sufficiently large gain. Limiting values are about 10% detuning relative to the Raman line width, and a fractional, unshifted Stokes intensity of about 1% .

Stable solitons will show temporal narrowing with a width inversely proportional to the square root of the total gain. For large gain the final width is independent of the initial width.

Unstable solitons will reach minimal width at a relative amplitude of .5, and begin to broaden on further propagation. The minimal width is proportional to the initial width and the initial relative amplitude.

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SOLITONS IN STIMULATED RAMAN SCATTERING: GENERATION AND CONTROL OF ULTRASHORT OPTICAL PULSES

FINAL REPORT

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U. S. ARMY RESEARCH OFFICE

CONTRACT NUMBER: DAAG29-85-K-0031

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FOREWORD

Some results obtained under this contract have not yet appeared in published form. It was decided to publish these together with additional theoretical work scheduled for the second phase of this project. For this reason the technical section B. of this report is some what longer, and kept in publication format.

TABLE OF CONTENTS

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A. RESEARCH OBJECTIVES

The research reported here deals with the problem of generation and propagation of Raman solitons in a medium with homogeneous broadening (coherence decay). The following questions are addressed:

1. Conditions under which a soliton excitation will develop in a Raman amplifier. Parameters are: amplitude and temporal profile of initial optical fields, Raman line width and detuning from Raman resonance.

2. Propagation of soliton excitations. Parameters are: soliton width and amplitude as a function of propagation distance and parameters under 1. above.

Methods employed are numerical solution of the transient Raman scattering equations, and analytical studies based on asymptotic perturbation theory.

B. TECHNICAL RESULTS

I. INTRODUCTION

Solitons in stimulated Raman scattering (SRS) were first observed experimentally in 1983 [1]. Stimulated Raman scattering involves at least two optical beams, one beam called the pump and a second beam called Stokes beam at a frequency lower than the pump. The difference in frequencies is equal or almost equal to the transition frequency of a Raman active medium, through which both beams propagate. If the medium is initially in the ground state ("Stokes scattering") photons are transferred from the pump beam to the Stokes beam as both beams propagate through the medium, leading to an amplification of the latter and to depletion of the pump (see [2] for a typical example). The medium absorbs the excess photon energy and is left in a partially excited state. The reverse process ("anti Stokes scattering") occurs if the medium is initially inverted.

Raman solitons are a coherent transient phenomenon in which photons are transferred back from the Stokes to the pump beam due to a pulse of coherent excitation of the medium. This excitation is itself sustained by the optical fields. The result is a stable localized non linear wave of excitation of both medium and optical fields, which can be observed experimentally as a pulse of pump radiation travelling through the medium in an envelope of Stokes radiation.

This phenomenon shows very close analogies to the phenomenon of self induced transparency (SIT) [3,4,5,6]. In fact an exact formal correspondence can be established between the equations describing the two phenomena [7,8,9,10]. The two levels of the atomic system in SIT correspond to the two levels of the photon system defined by the two beams in SRS, while the electric field in SIT corresponds to the coherent polarization in SRS. Also the roles of spatial and temporal variables is interchanged.

Similar phenomena have been investigated theoretically for coherent two photon propagation [10,11,12,13] and for three level systems, where all transitions are nearly resonant with an optical frequency ("simultons") [14,15].

Soliton solutions for SRS under various conditions and approximations have been studied for quite a while (see [8,9,10, 16,17,18] and papers quoted in these references). However it appears that no successful experimental studies have been undertaken until the recent discovery of these excitations [1] prompted further experimental work [19]. Additional theoretical studies have appeared since then, which addressed some of the observed interesting effects of coherence decay, like pulse narrowing [20,21,22,23] and the extension to higher order Stokes

generation and four wave mixing [24].

In section II we give the SRS equations used in this and other related work, discuss some of their relevant properties and distinguish different regimes by the time scale of variation of the physical quantities.

In section III we discuss the one soliton solutions, which are valid in the absence of coherence decay. We then give a simplified theoretical treatment of pulse narrowing in the presence of coherence decay and detuning, which is based on earlier work by Kaup [25].

In section IV we present numerical studies of soliton propagation for pulse widths comparable to or smaller than the coherence decay time, and compare these to the theoretical results of section III.

II. TRANSIENT SRS EQUATIONS AND THEIR PROPERTIES

The equations for transient SRS including coherence decay are given by [8,16]:

$$X_{\tau} = -\epsilon X + A_1 A_2^{*} , \tag{2.1}$$

$$A_{1\chi} = - X A_2 , \tag{2.2}$$

$$A_{2\chi} = X^{*} A_1 . \tag{2.3}$$

Here A_1 and A_2 are the Stokes and pump field and X is the off diagonal matrix element for the molecular Raman transition. τ and χ are time like and space like coordinates, which are related to time t and propagation distance z in the laboratory frame by $\tau = t - z/c$, $\chi = z$. Partial differentiation with respect to these coordinates is indicated by the corresponding subscript. The first term in (2.1) describes collisional coherence decay with decay time $1/\epsilon$, where ϵ is the angular Raman line width (HWHM Lorenzian) in radians per unit of time. Suitable units have been chosen to render all coupling constants equal to unity. We shall discuss below how to arrive at observable quantities from solutions of (2.1) to (2.3) by using intrinsic scales of time and length.

Certain effects are neglected in these equations. In addition to higher order Stokes generation and four wave mixing these are effects of dynamic Stark shift, medium polarization and medium saturation (population of the upper level). Their

validity is hence limited to low intensities, where these effects can be neglected. To compensate for this, the gain length can be increased to arrive at a desired value for the total gain. A formal expression for this is the fact that equations (2.1) to (2.3) are invariant under certain scaling transformations. These are:

$$\chi = \chi'/r \quad , \; X = r\, X' \; , \; A = \sqrt{r}\;\; A' \quad , \tag{2.4}$$

$$\tau = \tau'/s \quad , \; \epsilon = s\, \epsilon' \; , \; A = \sqrt{s}\;\; A' \quad . \tag{2.5}$$

Equation (2.4) can be considered as a change in length scale, or as a genuine invariance under simultaneous change of gain length and intensities. Equation (2.5) describes a change in time scale and becomes a physical invariance in the limiting case ϵ = 0 (hyper transient case, see below). Solutions have to be functions of quantities which are invariant under both transformations.

A convenient and physically intuitive way of checking this is a dimensional analysis. We denote the dimension of a quantity Q by [Q] and introduce independent temporal and spatial dimensions [τ] and [χ]. Then:

$$[X] = [\chi]^{-1} \; , \; [A_i\, A_j^{*}] = ([\chi][\tau])^{-1} \; , \; [\epsilon] = [\tau]^{-1} \tag{2.6}$$

Additional symmetries are:

$$A_1 = A_1' \;\; \exp[i\varphi(\tau) + i\,\varphi_0\,] \; , \quad A_2 = A_2' \exp[i\varphi(\tau)] , \tag{2.7}$$

$$X = X' \;\; \exp[i\,\varphi_0\,] \; :$$

$$\chi = -\chi' \; , \quad A_1 = A_2' \; , \quad X = X'^{*} \quad . \tag{2.8}$$

Equation (2.7) states that the fields are defined from (2.1) to (2.3) only up to an arbitrary time dependent phase factor, which is determined by the initial conditions. Equation (2.8) is a discrete symmetry which maps physical solutions into unphysical ones (pulses travelling at group velocity larger than the velocity c of light) and vice versa. Some solutions of (2.1) to (2.3) given in the literature [16,17] are actually unphysical, and can be transformed into physical ones by using (2.8).

We shall now discuss three different regimes distinguished by the time scale of variation for the physical quantities, in

which distinctly different types of phenomena occur:

Steady state: $X_\tau \ll \epsilon X$.

In this regime the matrix element is determined by the instantaneous value of the optical fields. The Stokes field sees gain, the pump field sees loss, and no soliton solutions are possible. The gain coeeficient g and the total linear gain constant G corresponding to propagation over distance χ at pump intensity I_0 are given below. G is dimension less, as required:

$$X = \epsilon^{-1} A_1 A_2^* , \qquad (2.1')$$

$$A_{1\chi} = - \epsilon^{-1} I_2 A_1 , \qquad (2.2')$$

$$A_{2\chi} = \epsilon \quad I_1 A_2 \ ; \quad I_j = A_j A_j^* \quad , \ j=1,2. \qquad (2.3')$$

$$g I_0 = 2 \epsilon' I_0 , \quad G = g I_0 \chi = 2 \epsilon' I_0 \chi \quad . \qquad (2.9)$$

Transient regime: $X_\tau \not\approx \epsilon X$.

No closed form analytical solutions are known except in the linear regime[22], where depletion of the pump is negligible and A_1 is assumed to be independent of χ . Numerical studies (see below) show however that for certain initial conditions solutions develop into the hyper transient regime discussed below (pulse narrowing), for which exact and approximate analytical methods of solution exist.

Hyper transient regime: $X_\tau \gg \epsilon X$.

In this regime the physical quantities vary on a time scale much shorter than the dephasing time $1/\epsilon$. In the limit $\epsilon = 0$ equations (2.1) to (2.3) are integrable, and exact solutions can be found by using the inverse spectral transform [16,17]. From these solutions the case of non vanishing ϵ can be studied through perturbation theoretical methods [25]. In the following we shall discuss the one soliton solution for the case $\epsilon = 0$ and its generalization to the case $\epsilon > 0$ by perturbation theory.

III. ONE SOLITON SOLUTIONS AND PULSE NARROWING IN THE HYPER TRANSIENT REGIME.

For $\epsilon = 0$ the SRS equations have soliton solutions [8,9,16, 17]. The one soliton solution is given by [16,17]:

$$X = \mu_I \exp(-iB) \operatorname{sech} A \,, \tag{3.1}$$

$$A = \sqrt{\mu_I \omega_R}\, \exp(iB) \operatorname{sech} A \,,$$

$$A = -\sqrt{\mu_I/\omega_R}\, [\, \omega_R \tanh A - i\, \omega_I \,] \,;$$

$$A = \omega_R \tau - \mu_I \chi + A_0 \,, \quad B = \omega_I \tau + \mu_R \chi \,, \tag{3.2}$$

$$\mu_I = I_0\, \omega_R / (\omega_R^2 + \omega_I^2) \,, \quad \mu_R = I_0\, \omega_I / (\omega_R^2 + \omega_I^2)$$

$$I_0 = A_1 A_1^* + A_2 A_2^* \,.$$

This solution describes a coherent excitation of both medium and fields travelling at a speed v smaller than the speed c of light. The total intensity (or photon density) in both fields is equal to I_0. The excitation can be observed as a localized pulse of pump radiation or as an (infinitely) extended pulse of Stokes radiation with a localized dip in intensity.

The temporal width of the excitation is equal to $\Delta\tau = 1/\omega_R$ which defines an intrinsic time scale. The solution will be valid if $\Delta\tau \ll 1/\epsilon$, where ϵ is the Raman line width. The frequency of the pump field is detuned from exact Raman resonance by $\Delta\omega = \omega_I$. At maximal height the pump pulse reaches a fraction ρ of total intensity given by $\rho = \omega_R^2 / (\omega_R^2 + \omega_I^2)$. Equations (3.3) summarize the observable parameters of the soliton and their relations. Note that we have introduced the Raman line width as an external time scale in order to express the speed of propagation v in terms of the experimental gain coefficient g (see(2.9)).

$$\Delta\tau = 1/\omega_R \,, \qquad \Delta\omega = \omega_I \,, \tag{3.3}$$

$$\rho = 1/(1 + \Delta\omega^2 \Delta\tau^2) \,,$$

$$1/v = 1/c + \mu_I/\omega_R = 1/c + g\, I_0\, \epsilon\, \Delta\tau^2 \rho/2 \,.$$

We now discuss the question of boundary conditions. We assume a Raman medium which is located in the half space $\chi > 0$ and is not excited initially. For optical pulses of finite duration the boundary conditions are:

$$X(\chi, \tau = 0) = 0 \,; \tag{3.4}$$

$$A_j(\chi = 0, \tau) = a_j(\tau) \,, \qquad \tau > 0 \,;$$

$$A_j(\chi = 0, \tau) = 0 \quad , \quad \tau \leq 0 .$$

Depending on the sign of A_0, the solution (3.1) will or will not satisfy these boundary conditions approximately. The two cases are:

$$A_0 > 0 \ , \quad A(\chi_0, \tau = 0) = 0 \quad \text{for} \quad \chi_0 > 0 . \tag{3.5}$$

In this case the solution will describe an initial excitation of the medium. For sufficiently large A_0 the fields at the entrance of the Raman cell ($\chi = 0$) will be almost equal to their asymptotic values. This situation corresponds to super radiant scattering from an initial macroscopic dipole moment, and has not been realized experimentally so far. It is not in agreement with (3.4).

$$A_0 < 0 \ , \quad A(\chi = 0, \tau_0) = 0 \quad \text{for} \quad \tau_0 > 0. \tag{3.6}$$

In this case the value for the polarization X in the medium ($\chi > 0$) will be exponentially small for sufficiently large absolute value of A_0. The fields in the Raman cell will be almost equal to their asymptotic values at $\tau = 0$, and assume values at soliton center ($A = 0$) for some finite $\tau > 0$. This solution is a close approximation to the initial conditions (3.5). Note that the field values can to a good approximation be set equal to zero in the region where solution (3.1) becomes asymptotic ($|A| \gg 1$). This then gives optical pulses of finite duration.

Solution (3.1) can be generalized to the case where $\epsilon > 0$ by using asymptotic perturbation theory in ϵ [21,23,25]. In view of the physical initial conditions (3.4) we would like to assume the fields and polarization as given and equal to their one soliton form at the entrance of the Raman cell ($\chi = 0$), and then calculate their evolution, as the pulses propagate into the cell ($\chi > 0$).

This problem however requires some non trivial reformulation [18,21,23]. An explicit discussion for the case $\omega_I = 0$ may be found in [23]. The reason for this is the fact that the additional term involving ϵ appears in the equation for X , which describes evolution in τ and not in χ . An alternative is to assume that all quantities are given and equal to their one soliton form through out the cell at $\tau = 0$, and then to calculate their time evolution. A formulation of asymptotic perturbation theory appropriate for this case is given in [25], and the results from this work can be readily applied to our

problem.

Although this approach does not solve the physical initial value problem, we argue that the results can still be applied. In fact the solutions are of one soliton form as functions of χ, however with more general τ -dependence of their coefficients. By linearizing these coefficients in τ about the soliton center given by $A = 0$, we obtain also one soliton solutions as functions of τ. The equation for the soliton trajectory, $A = 0$, is then solved to obtain the spatial position χ_0 of the soliton center as a function of the temporal position τ_0, and is used to express all coefficients obtained as functions of time. In detail this approach leads to the following results:

For $\epsilon > 0$ the one soliton contribution to the solution of (2.10 to (2.3) is of the form (3.1), where the argument functions A and B however have the more general form:

$$A = \alpha(\tau) - \mu_I(\tau)\chi, \quad B = \beta(\tau) + \mu_R(\tau)\chi \quad .(3.7)$$

The equations for the coefficients are [25]:

$$\mu_{I\tau} = -2\epsilon\mu_I \quad , \mu_{R\tau} = 0 \; , \qquad (3.8)$$

$$\alpha_\tau + 2\epsilon\alpha = \omega_R \quad , \quad \beta_\tau = \omega_I \quad . \qquad (3.9)$$

Equations (3.8) can be derived in a very straight forward way by using constants of motion [26]. Let us denote the coordinates of the soliton center by χ_0 and τ_0, and linearize the functions A and B about the soliton center:

$$A(\chi_0 + \chi', \tau_0 + \tau') = A_\chi \chi' + A_\tau \tau' \quad , \qquad (3.10)$$

$$B(\chi_0 + \chi', \tau_0 + \tau') = B_\chi \chi' + B_\tau \tau' \quad .$$

Then the coefficients of this expansion are obtained from (3.8) and (3.9) as:

$$A_\chi = -\mu_I(\tau_0) \; , \quad A_\tau = \omega_R(\tau_0) \; , \qquad (3.11)$$

$$B_\chi = \mu_R \quad , \quad B_\tau = \omega_I(\tau_0) \; .$$

It follows that the coefficients of the linearized arguments are exactly of the same form as for $\epsilon = 0$, except for the addi-

tional dependence on the temporal position τ_0 of the soliton. The solution for the soliton trajectory is:

$$\exp(4\epsilon\tau_0) = \exp(\alpha\chi_0) + (\mu_{I0}/\mu_R)^2[\exp(\alpha\chi_0) - 1] \qquad (3.12)$$

$$\approx 1 + \alpha'\chi_0 \quad , \text{ where}$$

$$\alpha = 4\epsilon\,\mu_R^2\, I_0^{-1} \quad , \quad \alpha' = 4\epsilon\,\mu_{I0}^2\, I_0^{-1} \quad .$$

The approximation given is valid for $\alpha\chi_0 \ll 1$ and agrees with the results obtained in [23]. In terms of observable parameters (3.3) the coefficients α and α' are given by:

$$\alpha = 2\rho_0 (1 - \rho_0)(\epsilon\Delta\tau)^2\, g\, I_0 \quad , \qquad (3.13)$$

$$\alpha' = 2\,(\rho_0\epsilon\Delta\tau)^2 g\, I_0 \quad .$$

The results for the relative pump intensity ρ at center, temporal soliton width $\Delta\tau$ and detuning $\Delta\omega$ are given below (initial values at = 0 are denoted by a subscript 0):

$$\rho = \rho_0 \exp(-\alpha\chi) \, , \qquad (3.14)$$

$$\Delta\tau^{-2} = \Delta\tau_0^{-2} \exp(-\alpha\chi) \,[1 + (\mu_{I0}/\mu_R)^2 (1 - \exp(-\alpha\chi))] , \qquad (3.15)$$

$$\Delta\omega = \Delta\omega_0 [1 + (\mu_{I0}/\mu_R)^2 (1 - \exp(-\alpha\chi))] \, . \qquad (3.16)$$

The maximal pump intensity decays exponentially with a rate that is approximately proportional to the residual Stokes intensity $(1 - \rho_0) I_0$, and hence to the square of the detuning. For zero detuning $\rho_0 = 1$ .This prediction is in excellent agreement with our numerical results (see below) and experimental observations at LANL [27]. The soliton width decreases initially and reaches a maximum when the relative pump intensity has decayed to $\rho = .5$. Its minimum is given by:

$$\Delta\tau^2_{MIN} = \Delta\tau_0^2\; 4\; \rho_0 (1 - \rho_0) \, . \qquad (3.17)$$

For $\rho_0 = 1$ or for zero detuning the soliton width will decrease with an inverse square root law:

$$\Delta\tau^2 = \Delta\tau_0^2 \,/[1 + \alpha'\chi] \, . \qquad (3.18)$$

In the limit of large gain for this case the width becomes independent of its initial value and is inversely proportional to the square root of the total gain:

$$\Delta\tau^2 = \epsilon^{-2} / (2 \rho_0^2 \ g I \chi) . \quad (3.19)$$

These results can be understood physically as the effects of a frequency chirp in the optical pulse, which is of dynamical origin. This is described by the explicit time dependence of the phase argument B above. The local frequency detuning is given by (3.16) above. Its maximum value for $\alpha\gamma \gg 1$ is:

$$\Delta\omega_{MAX} = \Delta\omega_D / (1 - \rho_0) . \quad (3.20)$$

As mentioned above these results are obtained by linearization and are valid if:

$$\epsilon \Delta\tau_0 \ll 1 \quad (3.21)$$

For the case of zero detuning they have been shown to be a good approximation to a more complete perturbation theory, which is valid to the extent that the solution can be described as a soliton with χ -dependent parameters, and distortions of the temporal pulses be neglected [23]. For the general case this question is difficult, and still under study.

IV. NUMERICAL STUDIES

Both soliton buildup from certain initial conditons and the propagation of fully developed solitons were studies extensively by numerical integration of equations (2.1) to (2.3). The features of soliton narrowing at increasing gain and soliton decay for detuning were verified and excellent agreement was found with the analytical results discussed in section III. Furthermore it was found that some analytical results obtained for propagation of solitons at large gain could also be applied to give information about soliton buildup and its dependence on the initial optical fields.

We start with a discussion of real solitons (no detuning). The initial Stokes field was assumed to undergo a change of sign in the form given by (3.1). The only dimensionless parameters are the threshold gain, or the quotient of input intensities and the width of the region of phase change, measured in units of

the dephasing time $1/\epsilon$. The resulting pulse shapes were found to be almost the same for a large region of variation of both parameters. The transition region, in which the Stokes intensity goes to zero, broadens in the linear regime considerably to about ten times the dephasing time. As soon as pump depletion sets in, however, pulse narrowing occurs. As the pump becomes depleted in the asymptotic temporal regions away from the transition region, a soliton develops with pulse shapes closely approximated by (3.1). The subsequent narrowing of this soliton at large gain is described very accurately by the approximate formulae (3.18) and (3.19).

A very detailed analysis and comparison between different analytical and numerical result si given in [23]. In figures (1.1) to (1.3) we show results from that analysis. Figure (1.1) shows a plot of the square of inverse soliton width versus gain. The uppermost broken curve is the straight line predicted from the approximate equation (3.18). The lower solid curve is from an extended perturbation theory discussed in [23]. The dotted line, which oscillates about the solid curve, gives numerical results. It is seen that the simple approximate theory gives very reasonable agreement.

Figure (1.2) shows results for the positon of the temporal soliton center, which depends only weakly (logarithmically) on propagation distance, rather than linear, as for $\epsilon = 0$. While this general dependence is confirmed, there is a slight discrepancy between the numerical results (upper dotted line), and the analytical results (lower broken and solid lines). This corresponds to a small constant error in soliton position, and is due to the difficulty of extracting the exact position of the one soliton contribution reliably from the total, distorted puls shapes.

Figure (1.3) shows the square norm of the pulse distortions as a function of gain. The initial pulse were assumed to be of exact soliton form, and distortions are seen to build up rapidly. Eventually however their amplitude decreases, and a pulse develops which is well approximated by the one soliton solution. This final situation always occurs if the soliton is allowed to build up from a weak Stokes pulse by Raman amplification.

In figures (1.4.a) to (1.4.d) we show numerical results for soliton buildup. Shown is the pump amplitude as a function of propagation distance. Gain values and amount of depletion are given below:

event	gain	pump depletion
a	8	3 %
b	12	56 %

c	16	99 %
d	20	100 %
d	30	100 %
d	40	100 %

The pulses in events a to c , although not soliton like, are seen to narrow with increasing gain [22]. For the three events under d the pulses approach the one soliton form, and their narrowing is in agreement with the analytical predictions discussed above.

For complex solitons (non zero detuning) soliton decay occurs in addition to initial narrowing. When the soliton has decayed to a relative maximum of ς = .5 , the width has a minimum and broadening begins. For our numerical studies we used quasi realistic pulse shapes. Three types of initial conditions were studied:

a) non zero detuning, real optical fields with change of sign for Stokes field by tanh function ("detuning");

b) zero detuning, real pump, complex Stokes field with constant imaginary part as for soliton solution ("leakage");

c) combination of a) and b) as in the one soliton solution, however at reduced total Stokes intensity.

Cases a) and b) give almost identical results for soliton decay, if the residual Stokes intensity (leakage) of $1 - \varsigma$ is equal to the square of the relative detuning $\Delta\omega / \epsilon$ (see 3.3). The name "leakage" is used since this case models an experimental situation in which the phase change is effected by a Pockels cell with residual transmission. Case c) gives a slight reduction in soliton decay.

In figure (2.1) we show superimposed pump intensities in five stages of propagation for case a) with a detuning of 1 % (relative to the Raman line width). The central pump pulse is seen to narrow initially, until is relative height reaches about .5 (fourth stage). In the final fifth stage no further reduction in width occurs (the apparent reduction is due to a shift in position and a decrease in amplitude).

In figure (2.2) we show sequences of propagation for a leakage of .1 % (2.2.a) and 1 % (2.2.b) (initial condition case b)). The first sequence is almost identical to the sequence in figure (2.1). The last sequence shows rapid soliton decay. Decay rates are in all cases in agreement with the theoretical prediction (3.14).

In figures (3.1) and (3.2) we compare events with increasing amount of detuning and leakage. Values are listed below:

event		a	b	c	d
(3.1)	detuning	1.5 %	3 %	30 %	300 %
(3.2)	leakage	0 %	.76 %	6.7 %	11.7 % .

Event (3.1.b) corresponds to a situation between events (3.2.a) and (3.2.b), while (3.1.c) lies between (3.2.c) and (3.2.d). Event (3.1.d) demonstrates a case of extreme detuning, in which no soliton formation is observed.

V. SUMMARY AND CONCLUSIONS

Soliton excitations in a Raman amplifier will develop, if a phase shift of nearly 180 degrees is introduced into the Stokes seed beam. In the absence of detuning and for an instantaneous phase shift of exactly 180 degrees the solitons will be stable. In the presence of detuning and for non instantaneous phase shifts of less than 180 degrees the solitons will be unstable and decay for sufficiently large gain. Limiting values are about 10% detuning relative to the Raman line width, and a fractional, unshifted Stokes intensity of about 1% .

Stable solitons will show temporal narrowing with a width inversely proportional to the square root of the total gain. For large gain the final width is independent of the initial width.

Unstable solitons will reach minimal width at a relative amplitude of .5, and begin to broaden on further propagation. The minimal width is proportional to the initial width and the initial relative amplitude.

These results show that media with large Raman line width will be important for practical applications. These will give both narrow final pulses, and pose less stringent requirements on detuning.

Figure (1.1)

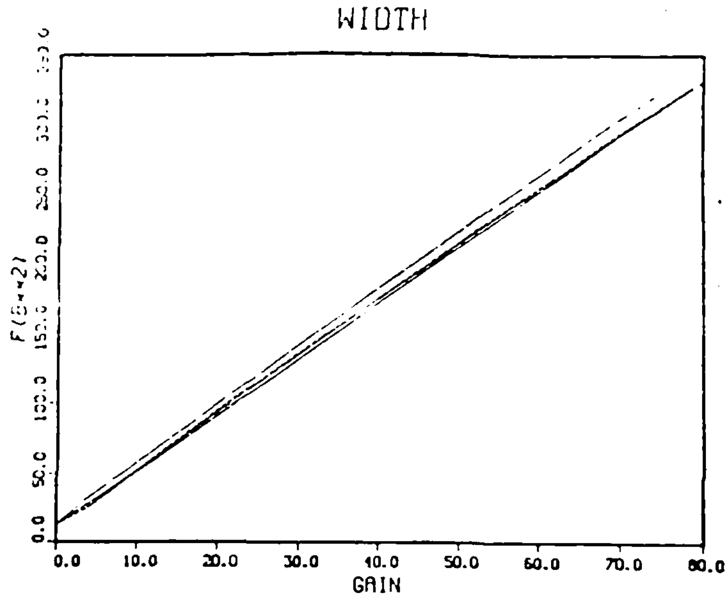


Figure (1.2)

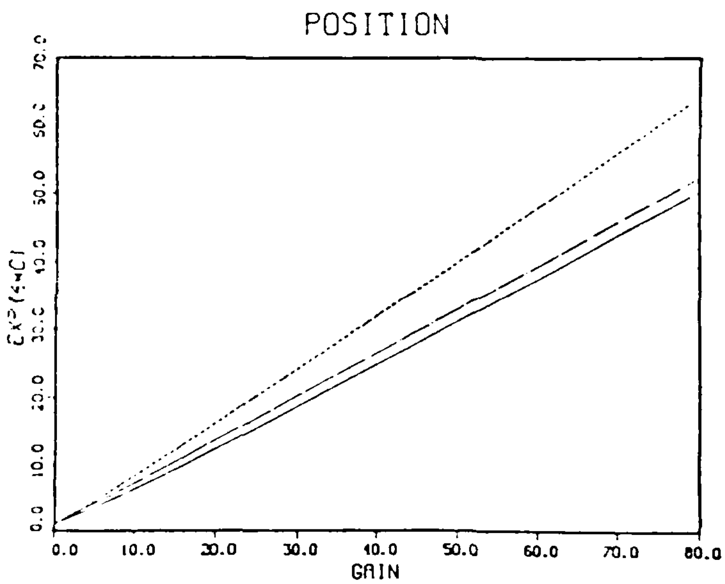
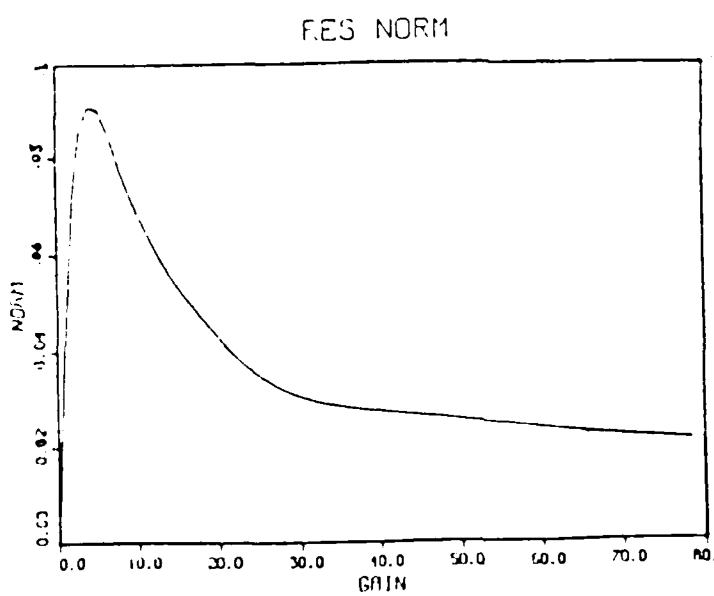


Figure (1.3)



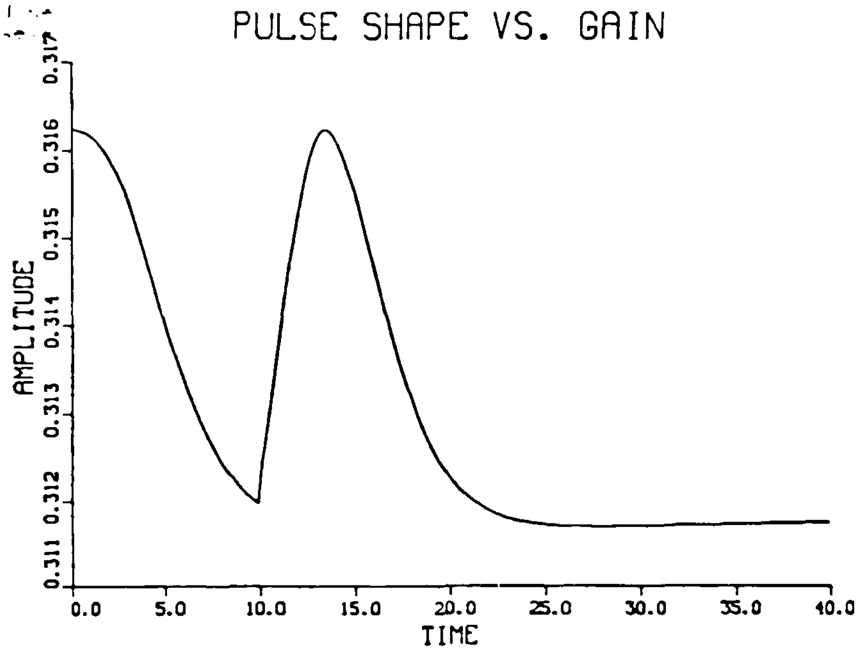


Figure (1.4.a)

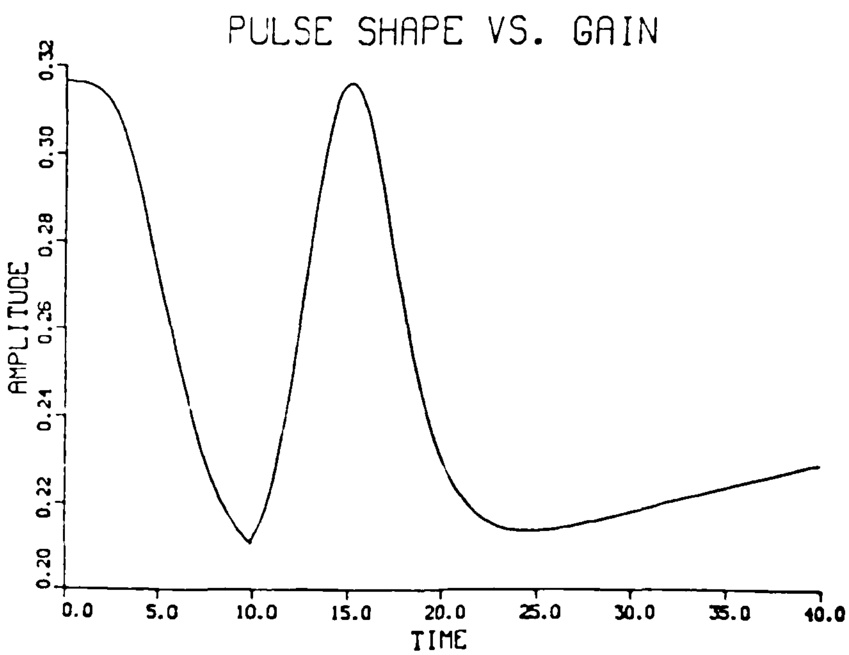


Figure (1.4.b)

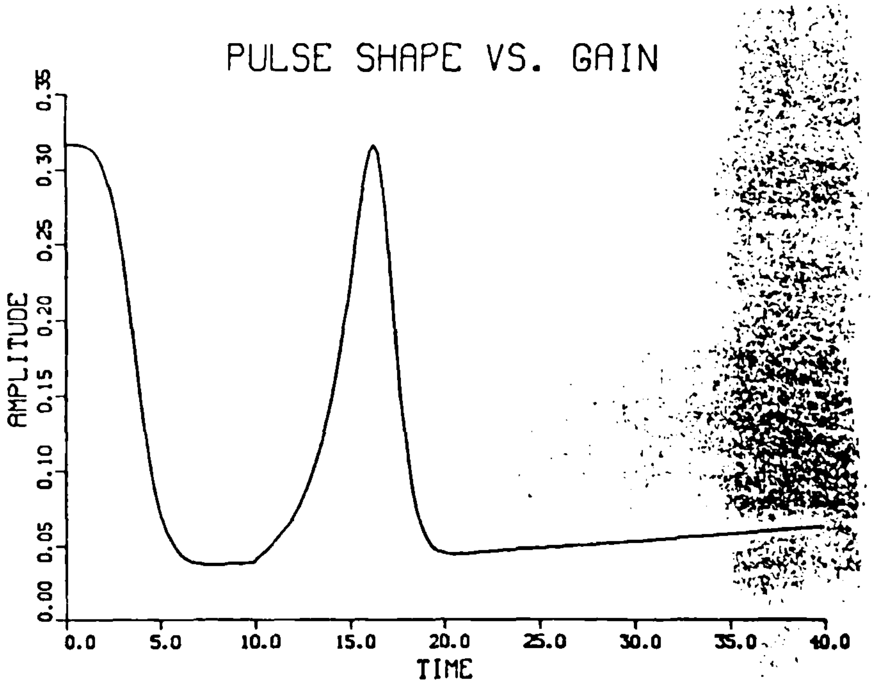


Figure (1.4.c)

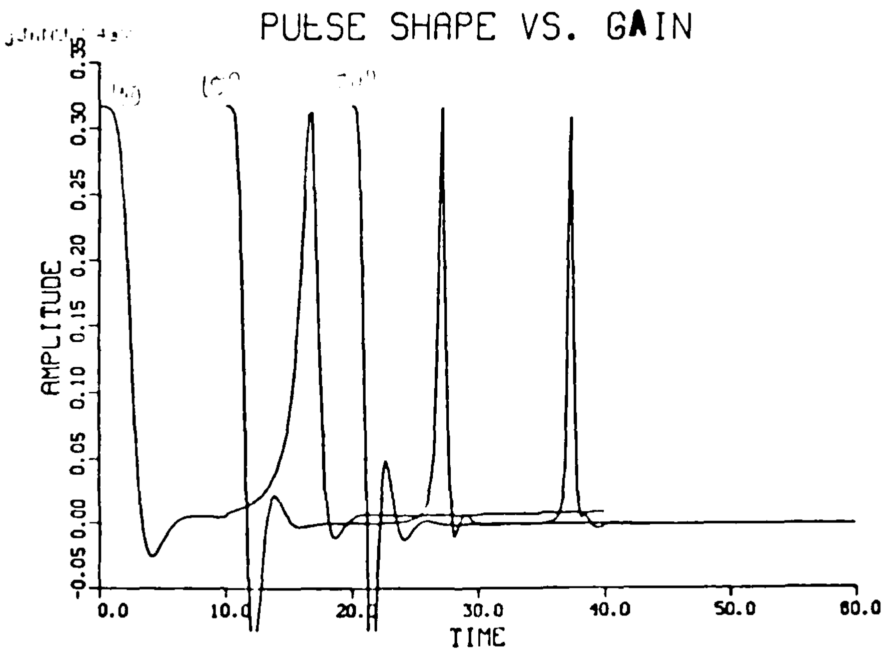


Figure (1.4.d)

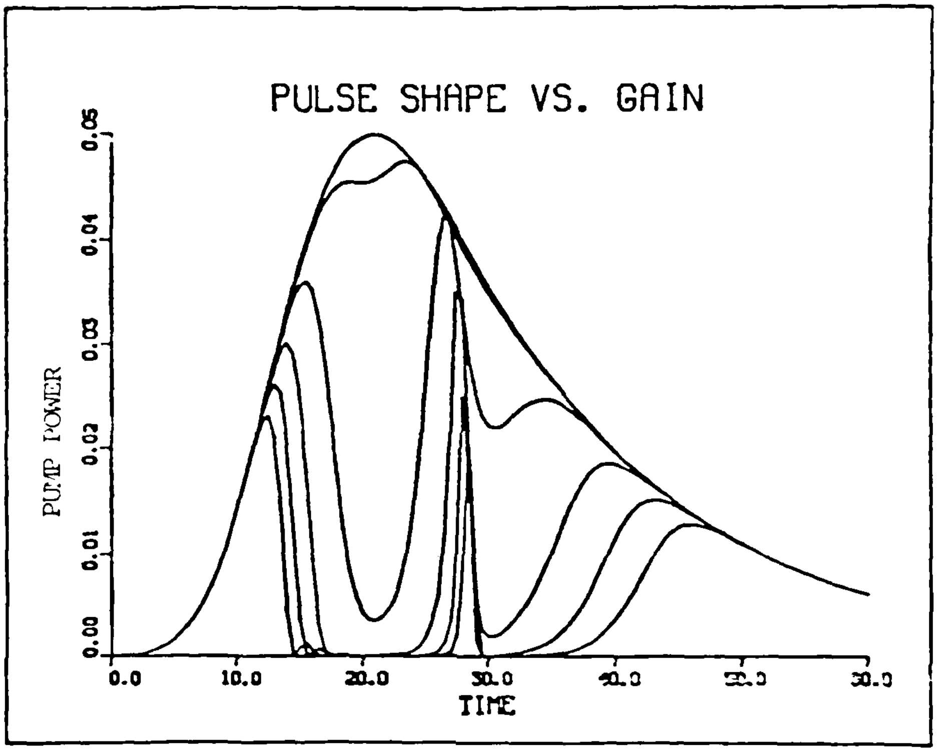


Figure (2.1)

Detuning = 1%.

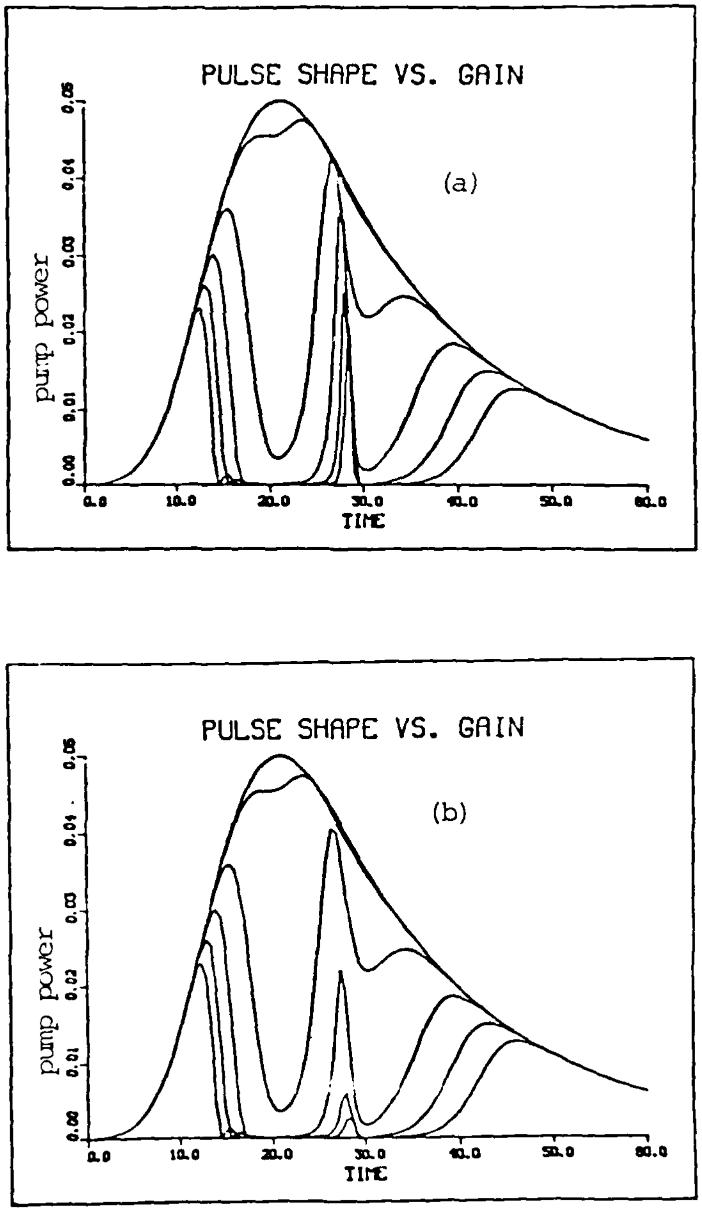


Figure (2.2)

(a) 0.1% leakage through the Pockel cell.

(b) 1.% leakage through the Pockels cell.

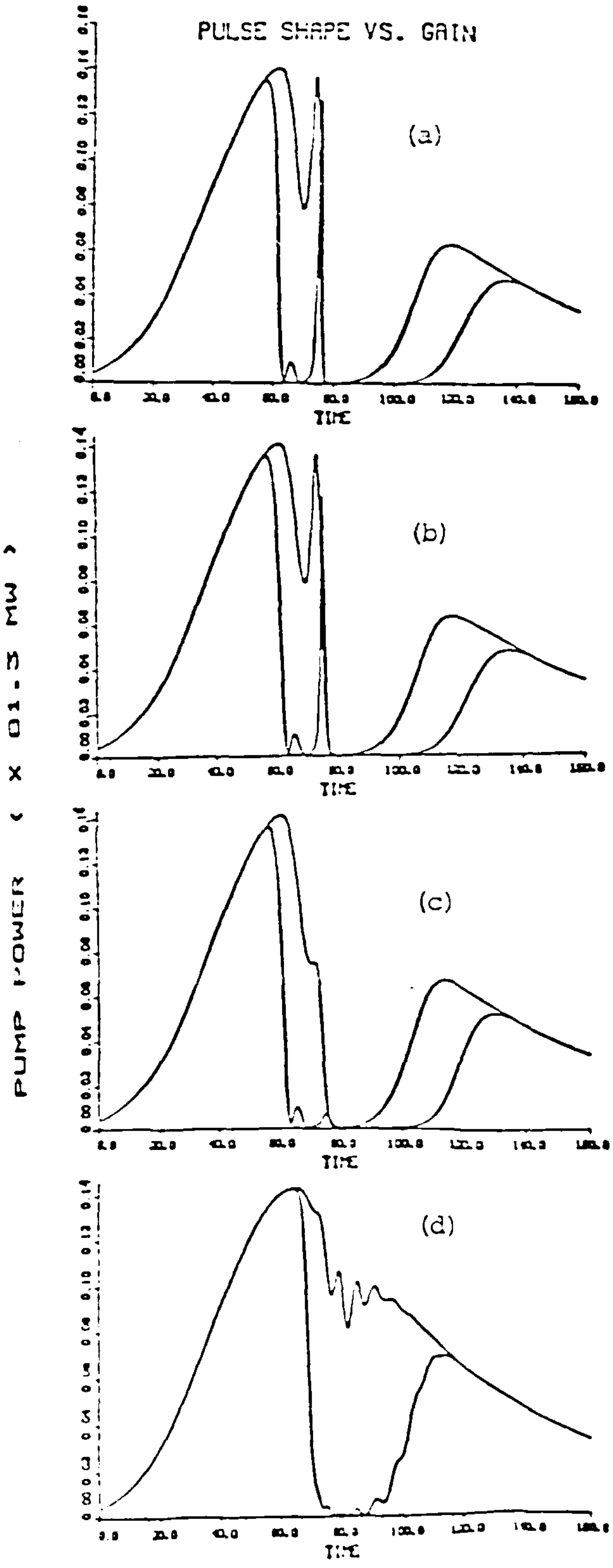


Figure (3.1)

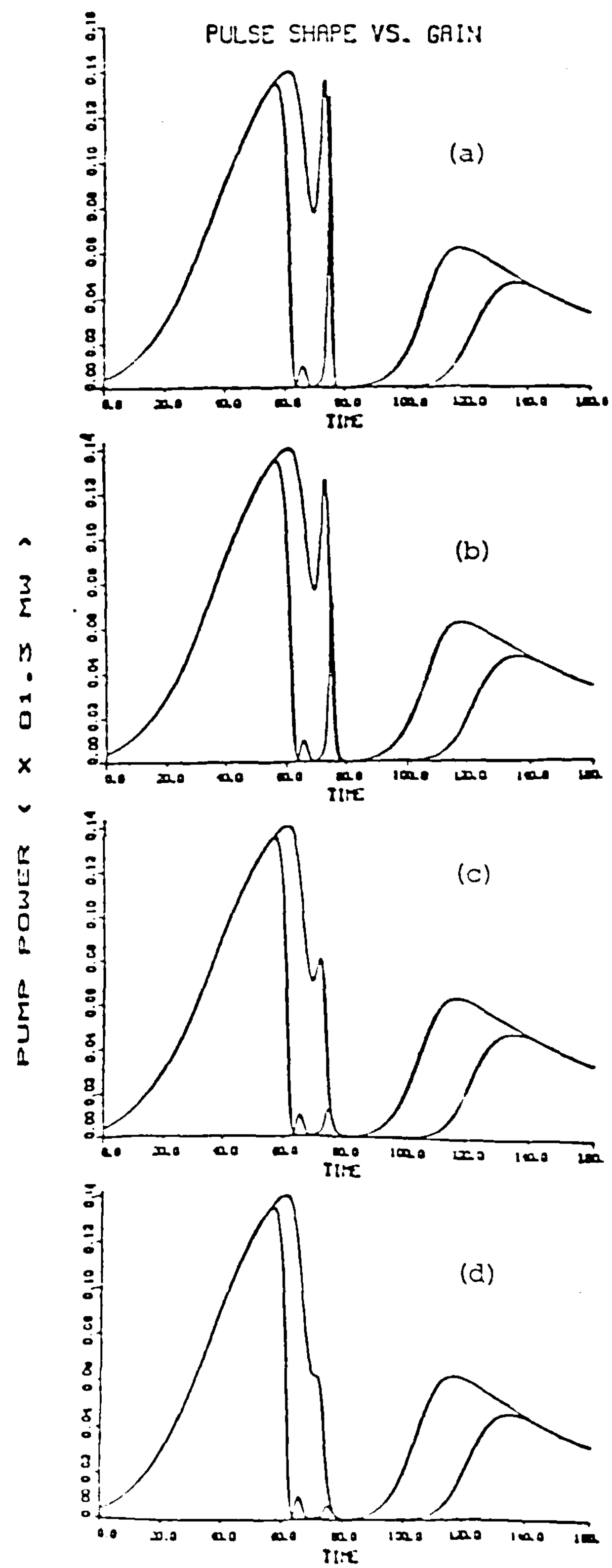


Figure (3.2)

C. PUBLICATIONS

K.Druhl, G.Alsing, "Effect of coherence relaxation on the propagation of optical solitons: An analytical and numerical case study on asymptotic perturbation theory", Physica D, May 1986 (in print).

D. LIST OF PERSONNEL

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